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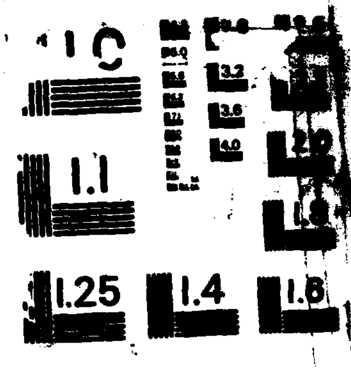

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Tapered Wiggler Analysis of High Gain Free Electron Laser Oscillators

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The development of lasers in which the active medium is a relativistic stream of free electrons has evoked much interest. The prospect of operating these continuously tunable, high gain devices at high efficiencies by appropriately tapering the wiggler field parameters has significantly increased the interest in these devices. Variable parameter wiggler free electron lasers have been previously considered for operation in the amplifier configuration and also in the micro-pulse, low gain operation in the oscillator configuration. The present analysis extends the range of application to the case of intense, long pulsed electron beams in which the self-consistently derived space charge potential (arising from the axial bunching of the particles in the beat wave) is not negligible. Also included in this analysis are arbitrarily relativistic beams (i.e., no ultra-relativistic beam approximation). We have obtained theoretical expressions for the nonlinear efficiency and present a comparison of these results with a fully nonlinear simulation.

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CONTENTS

Tapered Wiggler Analysis of High Gain Free Electron Laser Oscillators

I. Introduction

The development of lasers in which the active medium is a relativistic stream of free electrons has recently evoked much interest. The prospect of operating these continuously tunable, high gain devices at high efficiencies by appropriately tapering the wiggler field parameters has significantly increased the interest in these devices. The dramatic effect of efficiency enhancement schemes in free electron lasers relies on the maintenance of the co–resonance between the phase velocity of the ponderomotive wave with the decreasing velocity of the electron beam as it gives its energy to the radiation field. Variable parameter wiggler free electron lasers have been previously considered for operation in the amplifier configuration (Refs. 1-3) and also in the micro-pulse, low gain operation in the oscillator configuration (Refs. 4,5). The present analysis extends the range of application to the case of intense, long pulsed electron beams in which the self–consistently derived space charge potential (arising from the axial bunching of the particles in the beat wave) is not negligible. Also included in this analysis are arbitrarily relativistic beams (no ultra–relativistic beam approximation). Thusly, the simplifying assumptions that $\gamma \gg 1$ and $\beta_z \approx 1$ are not incorporated into the theory. The full relativistic effects of mildly relativistic beams, which are characteristic of beams from intense, long pulsed devices, are adequately considered. The analysis also includes the high gain operation in the oscillator configuration which is characterized by exponential field growth in the linear regime.

The organization of the paper is as follows. In Sec. II, we present the model and assumptions used in the numerical simulation. In Sec. III, we derive approximate theoretical expressions for the nonlinear efficiency. In Sec. IV, we present a summary of conclusions with specific examples from the simulation and comparison with theoretical expressions.

II. Simulation Model

We have conducted a fully relativistic, nonlinear simulation of the field and particle evolution in a variable parameter wiggler, high gain free electron laser oscillator. The simulation model which is kinetic in nature is coupled with Maxwell's equations to self-consistently include the evolution of the space charge waves. The wiggler field we consider

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is helically polarized in the plane transverse to the beam propagation direction, and is ideal in that the transverse variations in the vector potential are neglected.

$$\vec{A}_w(z) = -a_w(z)\left(\exp[-i \int_0^z k_w(z')dz'] \ \hat{e}_+ + c.c. \right), \tag{1}$$

where $a_w(z), k_w(z)$ are the arbitrarily prescribed wiggler amplitude and wave number respectively and also $\hat{e}_\pm = (\hat{e}_x \pm i\hat{e}_y)/2.$ and c.c. denotes complex conjugate. The radiation field model is also one-dimensional and is characterized by a temporally stationary amplitude with arbitrary spatial dependence.

$$\vec{A}_s(z,t) = a_s(z)\exp[i(k_s z - \omega_s t)] \ \hat{e}_+ + c.c. \tag{2}$$

The electrostatic space charge potential is also assumed to be temporally stationary with arbitrary spatial dependence.

$$\Phi(z,t) = \phi_1(z)\exp[i\{\int_0^z [k_s + k_w(z')]dz' - \omega_s t\}] \ + \ c.c. \tag{3}$$

The vector and scalar potentials of the electromagnetic wave evolve according to Maxwell's equations and the particles evolve in accordance with the kinetic distribution f (Ref. 6). Transverse effects are heuristically included by considering $\int dx \int dy \ f \ \propto \ \pi r_b^2$, $\int dx \int dy \ a_s \ \propto \ \pi w_0^2$ with r_b equal to the electron beam radius and $w_0 = (2z_R/k_s)^{1/2}$ is the spot size for an oscillator mode of Rayleigh range z_R. Inserting the driving currents and charge densities into Maxwell's equations and assuming i) a cold electron beam, ii) vector and scalar potentials that vary slowly as compared to an optical period and iii) small radiation fields, $a_s << a_w$, one obtains the following,

$$\frac{d\hat{a}_s}{dz} = \frac{i}{2}\left(\frac{\omega_b}{c}\right)^2 \frac{F\beta_{z0}}{k_s}\hat{a}_w(z)\left\langle \frac{e^{-i\psi}}{\gamma\beta_z} \right\rangle, \tag{4}$$

$$\left(\frac{d}{dz} + \frac{i(k_s + k_w)}{2}\right)\hat{\phi}_1 = -\frac{i}{2}\left(\frac{\omega_b}{c}\right)^2 \frac{\beta_{z0}}{(k_s + k_w)}\left\langle \frac{e^{-i\psi}}{\beta_z} \right\rangle, \tag{5}$$

where we have introduced the dimensionless potentials, $\hat{a}_s$, $\hat{a}_w$, $\hat{\phi}_1$ corresponding to $|e|(a_s, a_w, \phi_1)/mc^2$. The nonrelativistic beam plasma frequency is given by $\omega_b =$

$(4\pi ne^2/m)^{1/2}$ and $F = r_b^2/w_0^2$ is the filling factor. The integration over initial positions has been converted to an integral over initial phases, ψ_0, with $\langle\cdots\rangle = (1/2\pi)\int_{-\pi}^{\pi}\cdots d\psi_0$, and the phase is given by $\psi = \int_0^z [k_s + k_w(z') - \omega_s/v_z(z')]dz' + \psi_0$.

The evolution of the particle phase as determined by the Lorentz force equation is given by,

$$\psi'' = k_w' + \frac{\omega_s}{v_z^2}\frac{dv_z}{dz} = k_w' - \frac{k_s}{2\beta_z^3\gamma^2}\frac{d\hat{a}_w^2}{dz} - \frac{(1+\beta_z)k_s k_w}{\beta_z^3\gamma^2}\frac{\hat{a}_w}{2i}(\hat{a}_s e^{i\psi} - \hat{a}_s^* e^{-i\psi})$$
$$+ \frac{k_s(1-\beta_z^2)}{\beta_z^3\gamma}\left[e^{i\psi}(\frac{d}{dz} + i(k_s + k_w))\hat{\phi}_1 + e^{-i\psi}(\frac{d}{dz} - i(k_w + k_w))\hat{\phi}_1^*\right], \qquad (6)$$

where prime denotes differentiation with respect to z. The oscillator system of equations (Eqs. (4,5)) are solved with the following methods and boundary conditions. The ordinary differential equations for the fields and particles are solved with a four point Adams-Bashforth predictor corrector scheme which is initialized with a fourth order Runge-Kutta algorithm. After propagating an arbitrarily small radiation field from the entrance to the exit of the resonator, a portion of the wave (depending on the reflectivity of the cavity mirrors) is transmitted out of the system and the remaining wave is used to initialize the subsequent pass. In this manner we have been able to simulate the small signal exponential gain to the nonlinear saturated evolution of a tapered wiggler oscillator. Representative results are presented in Sec. IV.

III. Theoretical Analysis

The simulations are aided by monitoring a diagnostic quantity which may be obtained from Poynting's theorem. (Ref. 6)

$$\eta_R(z) \equiv \frac{P_R(z) - P_R(0)}{P_b} = \eta_E(z) \equiv \frac{\langle\gamma(0)\rangle - \langle\gamma(z)\rangle}{\gamma_0 - 1}, \qquad (7)$$

which expresses the fact that the energy loss by the particles must appear in the radiation field. These quantities are monitored at every axial grid of the simulation and provides an excellent check on the consistency of the field and particle evolution.

Defining the saturated state as the condition that the radiation power gained per pass equals the radiation power loss at the mirrors, one obtains

$$\frac{ck_s^2}{4\pi}\pi w_0^2[\mid a_s(L)\mid^2 - \mid a_s(0)\mid^2] = (1-R)\frac{ck_s^2}{4\pi}\pi w_0^2 \mid a_s(L)\mid^2$$

or

$$P_{Rsat}(L) = \frac{\eta_R(L)P_b}{(1-R)}, \tag{8}$$

where P_{Rsat} is the radiation power at saturation, P_b is the input electron beam power, η_R is the single pass radiation efficiency as defined in Poynting's theorem and R, is the power reflection coefficient for the cavity mirrors. This expression aptly indicates the possibility of achieving high radiation powers in high Q resonators with modest beam powers. The remaining theoretical sections are devoted to estimating, $\eta_R(L)$, the single pass radiation efficiency in terms of the prescribed beam and wiggler parameters.

The evolution of the radiation power can be obtained from the vector potential equation (4),

$$\frac{d\mid\hat{a}_s\mid}{dz} + i\phi_s'\mid\hat{a}_s\mid = \frac{i}{2}\left(\frac{\omega_b}{c}\right)^2\frac{\beta_{z0}F\hat{a}_w}{k_s}\left\langle\frac{e^{-i\psi_s}}{\beta_z\gamma}\right\rangle, \tag{9}$$

where ϕ_s is the radiation phase, $\hat{a}_s = \mid\hat{a}_s\mid e^{i\phi_s}$ and $\psi_s = \psi + \phi_s$. Multiplying the real part of Eq. (9) by $\mid\hat{a}_s\mid(c/4\pi)k_s^2\pi w_0^2$ one obtains

$$\frac{dP_R(z)}{dz} = \frac{P_b}{\gamma_0 - 1}k_s\hat{a}_w\mid\hat{a}_s\mid\left\langle\frac{\sin\psi_s}{\beta_z\gamma}\right\rangle. \tag{10}$$

III. a) Resonant Particle Approximation

The simplest estimate of the single pass efficiency as considered by Refs. 1,2 can be obtained by making the Compton and resonant phase approximations. The resonant phase approximation of Eq. (10) is given by,

$$\frac{dP_R(z)}{dz} = \frac{P_b}{\gamma_0 - 1}k_s\hat{a}_w\mid\hat{a}_s\mid\frac{\sin\psi_{sR}}{\beta_{z0}\gamma_R}. \tag{11}$$

The sine of the resonant phase is defined by evaluating the fixed point of the pendulum equation (Eq. (6)) at the resonant energy given by $\gamma_R^2 = k_s\mu^2/\beta_z(1+\beta_z)k_w$, with $\mu^2 = 1 + \hat{a}_w^2$.

$$0 = \phi_s'' + k_w' - \frac{k_s}{2\beta_z^3\gamma_R^2}\frac{d\hat{a}_w^2}{dz} - \frac{(1+\beta_z)k_w k_s}{\beta_z^3\gamma_R^2}\hat{a}_w \mid \hat{a}_s \mid \sin\psi_{sR}. \quad (12)$$

Neglecting the rate of change of the radiation phase as compared to the wiggler period and inserting the results from Eq. (12) into Eq. (11) yields,

$$\frac{dP_R(z)}{dz} = P_b\eta_0'(z), \quad (13)$$

where $\eta_0' \equiv \frac{\gamma_R(z)}{\gamma_0 - 1}\left[\frac{\beta_z^2 k_w'}{(1+\beta_z)k_w} - \frac{1}{2\mu^2}\frac{d\hat{a}_w^2}{dz}\right]$. Equation (13) can be integrated analytically and yields the following efficiency for the resonant particle approximation,

$$\eta_{Res} = \int_0^L \eta_0'(z)dz = \frac{P_R(L) - P_R(0)}{P_b} \approx \frac{2\beta_{z0}^2}{(1+\beta_{z0})}\frac{\gamma_R(0) - \gamma_R(L)}{\gamma_0 - 1}. \quad (14)$$

This efficiency is the relativistic generalization of the bucket efficiency of Refs. 1,2 and yields efficiencies which are of the order of 70 percent larger than the efficiencies obtained from the simulation of continuous electron beams. Although, when the simulations are performed for prebunched beams characterized by macro-particles injected at the resonant phase, Eq. (14) yields very good agreement with that simulation efficiency.

III. b) Deeply Trapped Particle Compton Limit

In order to better approximate the ensemble average over initial phases, the assumption of a macro–particle at the resonant phase will be replaced with an explicit average over the orbits of deeply trapped particles. As a first example we shall consider again the Compton approximation. Expanding the pendulum equation, Eq. (6), about the resonant phase, $\psi_s = \psi_{sR} + \delta\psi$ at the resonant energy $\gamma_R^2 = \mu^2 k_s/\beta_z(1+\beta_z)k_w$ yields,

$$\delta\psi'' = k_w' - \frac{k_s}{2\beta_z^3\gamma_R^2}\frac{d\hat{a}_w^2}{dz} - \psi_{sR}''(z) - \frac{(1+\beta_z)k_w k_s}{\beta_z^3\gamma_R^2}\hat{a}_w \mid \hat{a}_s \mid \{\sin\psi_{sR} + \cos\psi_{sR}\ \delta\psi\}, \quad (15)$$

where we shall explicitly retain the drift and acceleration of the resonant phase, $\psi'_{sR} \neq 0, \psi''_{sR} \neq 0$. Using the definition of the sine of the resonant phase, Eq. (12), the pendulum equation becomes,

$$\delta\psi'' = -\psi''_{sR}(z) - k^2_{syn}(z)\delta\psi, \tag{16}$$

where $k^2_{syn}(z) \equiv (1+\beta_z)k_w k_s \hat{a}_w \mid \hat{a}_s \mid \cos\psi_{sR}/\beta_z^3\gamma_R^2$. The oscillation of the deeply trapped particle can be explicitly solved by multiple scale techniques which exploit the desperate time scales of the synchrotron bouncing of the particle and the adiabatic variation of wiggler parameters. The solution is given by,

$$\delta\psi = -\frac{\psi''_{sR}(z)}{k^2_{syn}(z)} + \xi(\psi_0, z)\cos\int_0^z k_{syn}(z')dz' - \rho(z)\sin\int_0^z k_{syn}(z')dz' \tag{17}$$

where $\xi(\psi_0, z) = \sqrt{k_{syn}(0)}(\psi_0 - \psi_{sR}(0) + \psi''_{sR}(0)/k^2_{syn}(0))/\sqrt{k_{syn}(z)}$, $\rho(z) = \psi'_{sR}(0)/\sqrt{k_{syn}(0)k_{syn}(z)}$ and we have chosen boundary conditions corresponding to $\psi_s(0) = \psi_0$ and $\psi'_s(0) = 0$. Making use of this solution for the ensemble average yields,

$$\begin{aligned}\langle e^{-i\psi_s}\rangle =& e^{-i(\psi_{sR}(z) - \psi''_{sR}(z)/k^2_{syn}(z))}\int_{-\pi}^{\pi}\frac{d\psi_0}{2\pi}\sum_{k=-\infty}^{\infty}(i)^k J_k(\xi(\psi_0, z))\\ & e^{-ik\int_0^z k_{syn}(z')dz'}\sum_{\ell=-\infty}^{\infty} e^{i\ell\int_0^z k_{syn}(z')dz'} J_\ell(\rho(z)),\end{aligned} \tag{18}$$

where $J_m(x)$ is the Bessel function of the first kind of order m. By neglecting the oscillations at the sideband harmonics one obtains,

$$\langle\sin\psi_s\rangle = \sin(\psi_{sR}(z) - \psi''_{sR}(z)/k^2_{syn}(z))J_0(\rho(z))\Im, \tag{19}$$

$$\text{with} \qquad \Im \equiv \int_{-\pi}^{\pi}\frac{d\psi_0}{2\pi}J_0(\xi(\psi_0, z)),$$

$$\text{and} \qquad \int_0^\lambda d\lambda' J_0(\lambda') = \lambda J_0(\lambda) + \frac{\pi\lambda}{2}\left[H_0(\lambda)J_1(\lambda) - H_1(\lambda)J_0(\lambda)\right],$$

where $H_k(x)$ is the Struve function of order k. The results of the ensemble average, when included in the power evolution equation, yields,

$$\frac{dP_R(z)}{dz} = \eta'_I P_b - \eta'_{II}\sqrt{P_b P_R(z)}, \tag{20}$$

$$\text{where} \qquad \eta'_I \equiv \Im J_0\left(\rho(z)\right)\cos\frac{\psi''_{sR}(z)}{k^2_{syn}(z)}\eta'_0(z)$$

$$\text{and} \qquad \eta'_{II} \equiv \frac{\gamma_R}{\gamma_0 - 1}\Im J_0\left(\rho(z)\right)\sin\frac{\psi''_{sR}(z)}{k^2_{syn}(z)}\cos\psi_{sR}(z)\frac{\hat{a}_w\omega_b}{c\gamma_R^2}\left(\frac{\gamma_0 - 1}{\beta_z F}\right)^{1/2}.$$

Although Eq. (20) can be solved analytically for the case of constant η'_I and η'_{II}, the transcendental solution for $P_R(z)$ makes it impossible to evaluate the single pass efficiency, $\eta_R(L) = (P_R(0) - P_R(L))/P_b$. For this reason we choose to retain the slow variation of η'_I and η'_{II} and solve Eq. (20) numerically and compare the results to the simulations. This efficiency yields results that are within 10 percent of the simulation efficiency for the Compton regime operation. For parameters of physical interest $\eta'_I \gg \eta'_{II}$ and the evolution of Eq. (20) is dominated by the first term. The deeply trapped particle approximation to the efficiency is dominated by the bucket efficiency times the product of three form factors, $\Im$, $J_0(\rho(z))$and $\cos\psi''_{sR}(z)/k^2_{syn}(z)$, which are due to the average motion in the bucket, drift and acceleration of the bucket respectively.

III. c) Deeply Trapped Particle Raman Limit

The space charge effects can be easily included with this technique of deeply trapped orbit averaging. The pendulum equation in the presence of space charge waves and expanded about the resonant phase and energy is given by,

$$\begin{aligned}\delta\psi'' = &-\psi''_{sR} + k'_w - \frac{k_s}{2\beta_z^3\gamma_R^2}\frac{d\hat{a}_w^2}{dz} - \frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R}(k_s + k_w)|\phi_1|\sin(\psi_{sR} + \varphi - \phi_s)\\ &- \frac{(1+\beta_z)k_w k_s}{\beta_z^3\gamma_R^2}\hat{a}_w|\hat{a}_s|\sin\psi_{sR} - \left[\frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R}(k_s + k_w)|\phi_1|\cos(\psi_{sR} + \varphi - \phi_s)\right.\\ &\left.+ \frac{(1+\beta_z)k_w k_s}{\beta_z^3\gamma_R^2}\hat{a}_w|\hat{a}_s|\cos\psi_{sR}\right]\delta\psi.\end{aligned} \tag{21}$$

where φ is the phase of the electrostatic potential, $\hat{\phi}_1 = |\hat{\phi}_1|e^{i\varphi}$, and we have neglected the rate of change of the space charge potential as compared to the ponderomotive wavelength. By making use of the resonant phase definition in the Compton regime, $k'_w - k_s/(2\beta_z^3\gamma_R^2)d\hat{a}_w^2/dz = (1+\beta_z)k_w k_s \hat{a}_w|\hat{a}_s|\sin\psi_{sR}/\beta_z^3\gamma_R^2$, the pendulum equation can be cast in the form,

$$\delta\psi'' = \alpha(z) - \tilde{k}_{syn}^2\delta\psi, \tag{22}$$

where $$\alpha(z) \equiv -\psi''_{sR} - \frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R^2}(k_s+k_w)|\hat{\phi}_1|\sin(\psi_{sR}+\varphi-\phi_s),$$

$$\tilde{k}_{syn}^2(z) = \frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R^2}(k_s+k_w)|\hat{\phi}_1|\cos(\psi_{sR}+\varphi-\phi_s) + \frac{(1+\beta_z)k_w k_s}{\beta_z^3\gamma_R^2}\hat{a}_w|\hat{a}_s|\cos\psi_{sR}.$$

The magnitude of the electrostatic potential, $|\hat{\phi}_1|$, and the phase shift between the electrostatic and radiation phases, $\varphi - \phi_s$ can be estimated from the evolution equation for the space charge potential.

$$|\hat{\phi}_1| = \frac{-\omega_b^2}{(k_s+k_w)^2c^2}\langle\cos(\psi_s+\varphi-\phi_s) - i\sin(\psi_s+\varphi-\phi_s)\rangle. \tag{23}$$

Since the left-hand side of Eq. (23) is real, one has that $\langle\sin(\psi_s+\varphi-\phi_s)\rangle = 0$. Using the deeply trapped orbits of the Compton approximation to evaluate the ensemble average yields,

$$\langle\sin(\psi_s+\varphi-\phi_s)\rangle = 3J_0(\rho(z))\sin(\psi_{sR}+\varphi-\phi_s-\psi''_{sR}/k_{syn}^2) = 0, \tag{24}$$

$$\psi_{sR}+\varphi-\phi_s - \frac{\psi''_{sR}}{k_{syn}^2} = \pi.$$

Thusly, the coefficients in the space charge pendulum equation are given by,

$$\alpha(z) = -\psi''_{sR} + \frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R}\frac{\omega_b^2}{(k_s+k_w)c^2}3J_0(\rho(z))\sin\frac{\psi''_{sR}(z)}{k_{syn}^2(z)}, \tag{25}$$

$$\tilde{k}_{syn}^2(z) = -\frac{2k_s(1-\beta_z^2)}{\beta_z^3\gamma_R}\frac{\omega_b^2}{(k_s+k_w)c^2}\Im J_0\left(\rho(z)\right)\cos\frac{\psi_{sR}''(z)}{\tilde{k}_{syn}^2(z)}$$
$$+\frac{(1+\beta_z)k_sk_w}{\beta_z^3\gamma_R^2}\hat{a}_w|\hat{a}_s|\cos\psi_{sR}, \tag{26}$$

where the space charge potential modifies the synchrotron wave vector $\tilde{k}_{syn}$ by increasing the synchrotron period due to electrostatic repulsion, and also increases the effective bucket acceleration $\alpha(z)$. Again making use of multiple scale techniques the solution to the pendulum equation is given by,

$$\delta\psi = \frac{\alpha(z)}{\tilde{k}_{syn}^2(z)} + \tilde{\xi}(\psi_0, z)\cos\int_0^z \tilde{k}_{syn}(z')dz' - \tilde{\rho}(z)\sin\int_0^z \tilde{k}_{syn}(z')dz', \tag{27}$$

where $\tilde{\xi}(\psi_0, z) = \sqrt{\tilde{k}_{syn}(0)}(\psi_0 - \psi_{sR}(0) - \alpha(0)/\tilde{k}_{syn}^2(0))/\sqrt{\tilde{k}_{syn}(z)}$ and $\tilde{\rho}(z) = \psi_{sR}'(0)/\sqrt{\tilde{k}_{syn}(0)\tilde{k}_{syn}(z)}$. Performing the ensemble average with these orbits, the power evolution equation becomes,

$$\frac{dP_R(z)}{dz} = {}_{sp}\eta_I' P_b + {}_{sp}\eta_{II}'\sqrt{P_bP_R(z)}, \tag{28}$$

where $${}_{sp}\eta_I' = \tilde{\Im}J_0\left(\tilde{\rho}(z)\right)\cos\frac{\alpha(z)}{\tilde{k}_{syn}^2(z)}\eta_0'(z)$$

and $${}_{sp}\eta_{II}' = \frac{\gamma_R}{\gamma-1}\tilde{\Im}J_0\left(\tilde{\rho}(z)\right)\sin\frac{\alpha(z)}{\tilde{k}_{syn}^2(z)}\cos\psi_{sR}\ \frac{\hat{a}_w\omega_b}{c\gamma_R^2}\left(\frac{\gamma_0-1}{\beta_zF}\right)^{1/2}$$

and $$\tilde{\Im} = \int_{-\pi}^{\pi}\frac{d\psi_0}{2\pi}J_0\left(\tilde{\xi}(\psi_0, z)\right).$$

The numerical solution to Eq. (28) is compared to the simulation results in the space charge dominated regime. The comparison presented in Sec. IV is characterized by agreement between theory and simulation on the order of 10 percent.

IV. Results and Summary

The simulation model is capable of tracking the evolution of the radiation field from the small signal exponential gain through the nonlinear saturated state of the oscillator.

During the small signal evolution of the fields for which the growth rate is constant, the growth rate from the simulation is compared to the growth rate obtained from the solution of the linear dispersion relation. This comparison is presented in Fig. 1 for the case of, beam current $I = 300A$, beam energy $\gamma = 2.5$, wiggler field strength $B_w = 1000G$, wiggler period $\ell_w = 4.0cm$, filling factor $F = 0.2$ and no taper of the wiggler amplitude or period. The comparison of the simulation and theory yields excellent agreement throughout the unstable spectrum.

An example of the small signal to nonlinear evolution of the fields is presented in Fig. 2. In this example the field evolves over 200 radiation bounce times for the case of beam current $I = 300A$, beam energy $\gamma = 2.5$, wiggler field strength $B_w = 1000G$, initial wiggler period $\ell_w = 4.0cm$, filling factor $F = 0.2$, linear taper of the wiggler period of 50 percent and mirror reflectivity $R = 0.98$. The saturated power and efficiency are $1.9GW$ and 17 percent respectively; and exhibits the scaling indicated in Eq. (28).

The single pass efficiency exhibits a weak dependence on the radiation power at the entrance of the resonator. In Fig. 3 we present a comparison of the single pass efficiency from the simulations and theory as a function of the input power. As noted in Sec. III. b), the theoretical evaluation of the efficiency neglects oscillations at the harmonics of the synchrotron period. The simulation model contains the complete dynamics of the electrons as dictated by the pendulum equation and the simulation efficiency thusly exhibits oscillations at synchrotron harmonics with an amplitude of the order of the well depth. The bars on the data points from the simulation indicate the amplitude of this oscillation in the efficiency. The resulting comparison indicates agreement on the order of 10 percent between the simulation and theoretical expressions.

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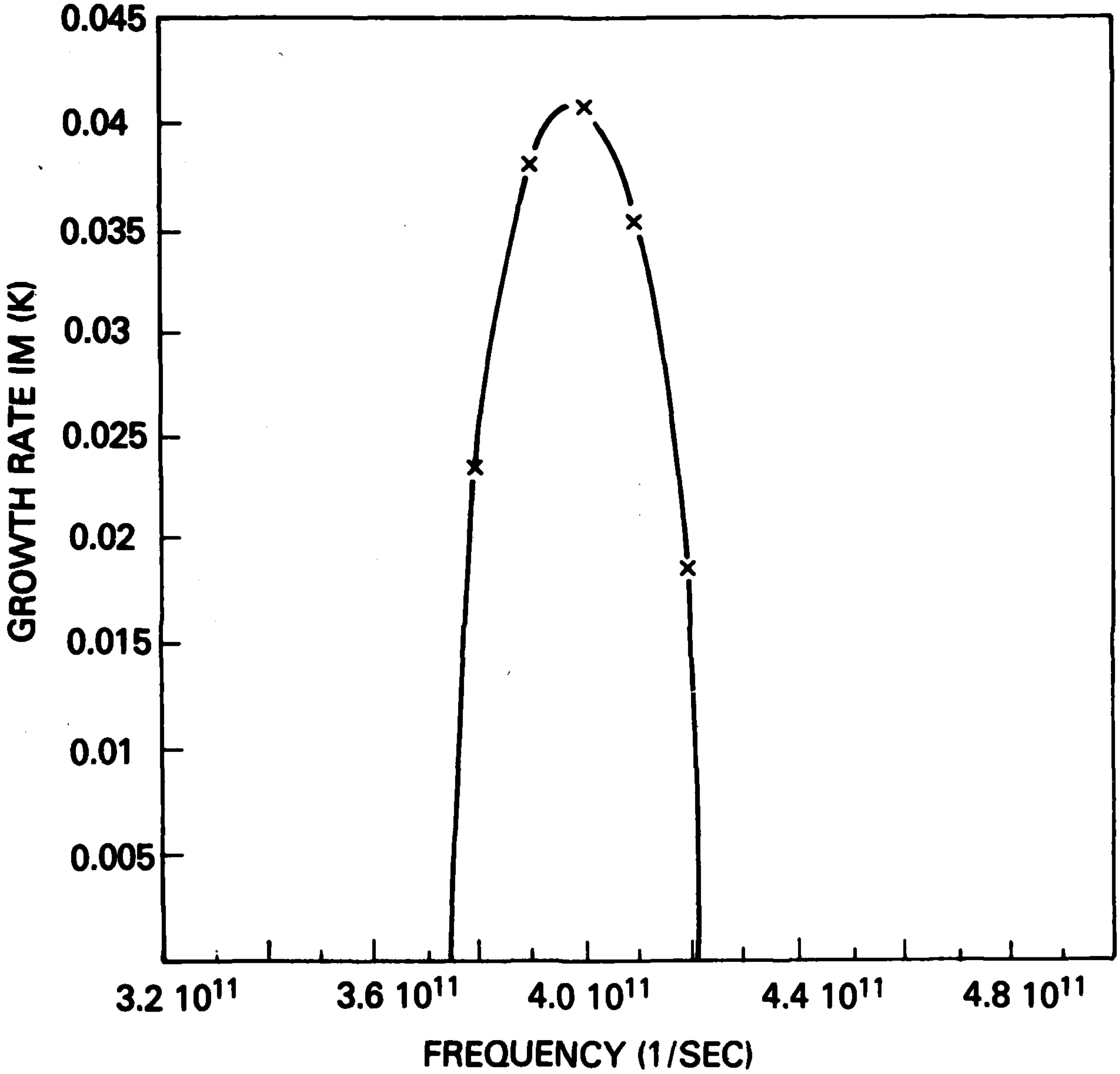

Figure 1. Comparison of theoretical and simulation growth rates for the case of $I = 300A$, $\gamma = 2.5$, $B_w = 1.0kG$, $\ell_w = 4.0cm$ and $F = 0.2$.

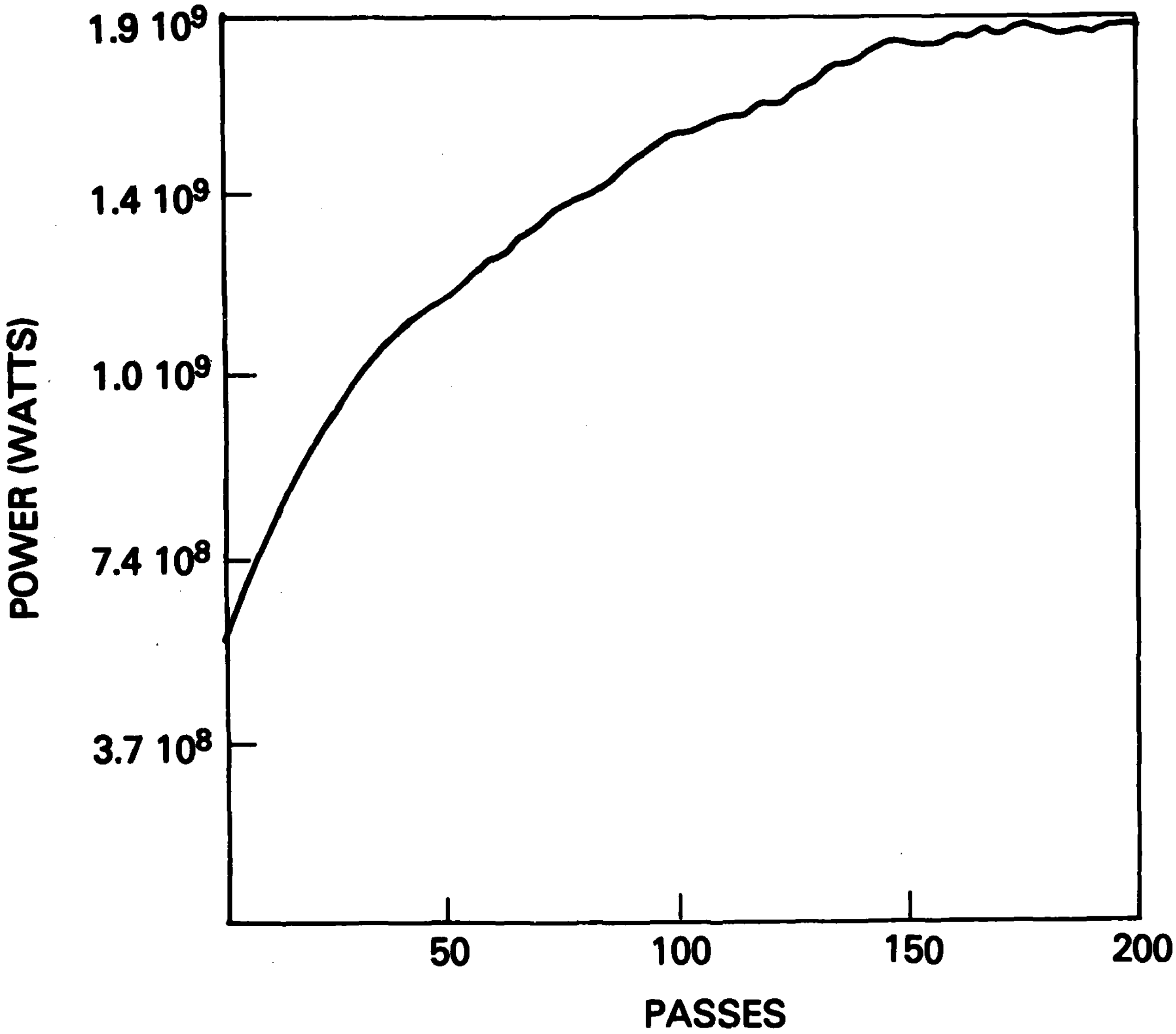

Figure 2. Radiation field evolution for the case of $I = 300A$, $\gamma = 2.5$, $B_w = 1.0kG$, $\ell_w = 4.0cm$, $F = 0.2$, $R = 0.98$ and 50% taper.

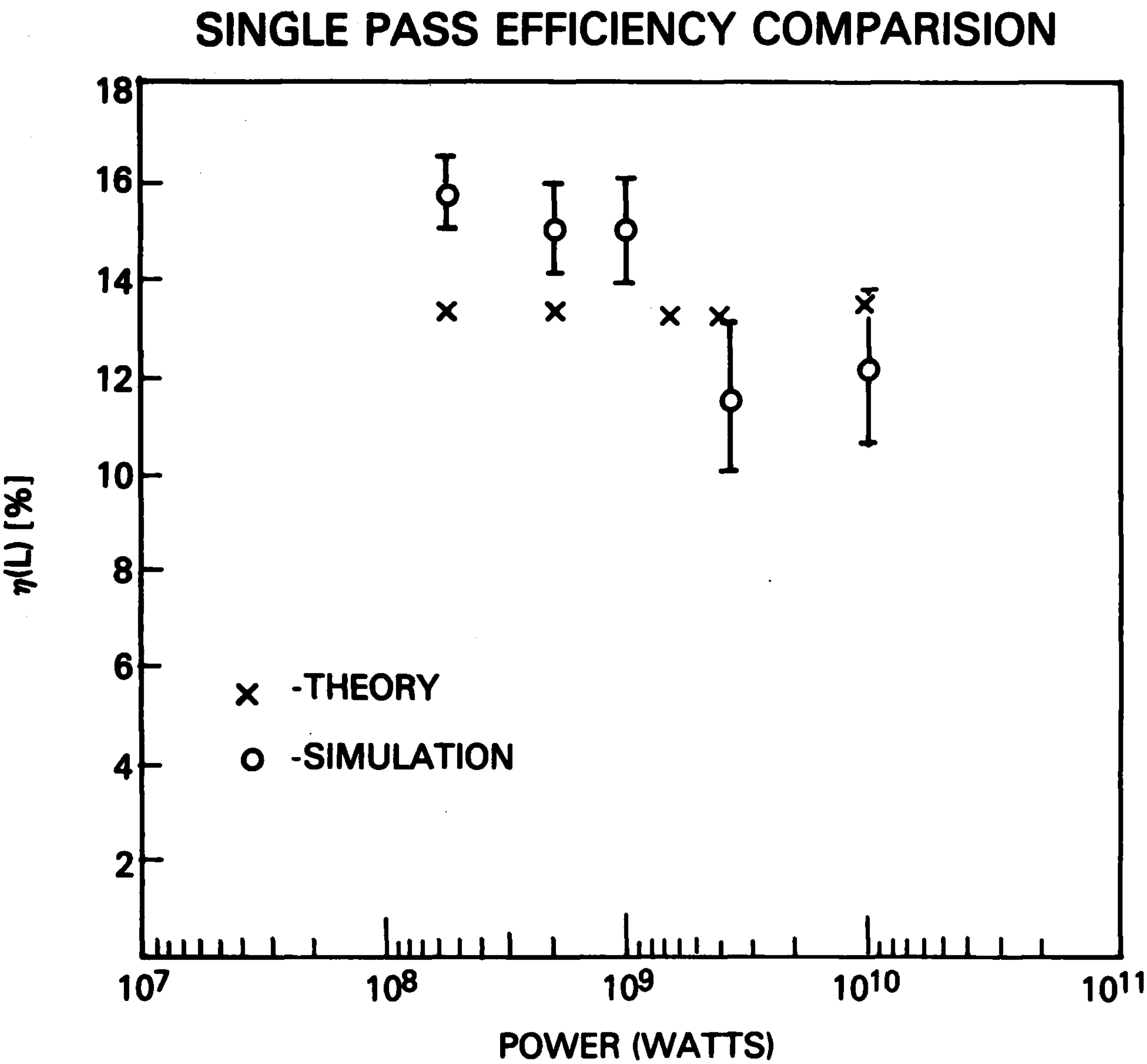

Figure 3. Comparison of theoretical and simulation values for the single pass efficiency as a function of input power.

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